

Brane solutions in strings with broken supersymmetry and dilaton tadpoles

E. Dudas^a and J. Mourad^b

^a *LPT[†], Bât. 210, Univ. de Paris-Sud, F-91405 Orsay, France*

^b *LPTM, Site Neuville III, Univ. de Cergy-Pontoise, Neuville sur Oise
F-95031 Cergy-Pontoise, France*

Abstract

The tachyon-free nonsupersymmetric string theories in ten dimensions have dilaton tadpoles which forbid a Minkowski vacuum. We determine the maximally symmetric backgrounds for the $USp(32)$ Type I string and the $SO(16) \times SO(16)$ heterotic string. The static solutions exhibit nine dimensional Poincaré symmetry and have finite 9D Planck and Yang-Mills constants. The low energy geometry is given by a ten dimensional manifold with two boundaries separated by a finite distance which suggests a spontaneous compactification of the ten dimensional string theory.

[†]Unité mixte de recherche du CNRS (UMR 8627).

1 Introduction

Nonsupersymmetric string models are generically plagued by divergences which raise the question of their quantum consistency. In particular, ten dimensional nonsupersymmetric models have all dilaton tadpoles and some of them have tachyons in the spectrum. It is however believed that some tadpoles do not signal an internal inconsistency of the theory but merely a background redefinition [1]. In particular, in orientifolds of Type II theories there should be a difference between tadpoles of Ramond-Ramond (RR) closed fields and tadpoles of Neveu-Schwarz Neveu-Schwarz (NS-NS) closed fields. While the first ones cannot be cured by a background redefinition and signal an internal inconsistency of the theory asking therefore always to be cancelled [2], the latter ones could in principle be cured by a background redefinition. The NS-NS tadpoles remove flat directions, generate potentials for the corresponding fields (for example dilaton in ten dimensions) and break supersymmetry. The difference between RR and NS-NS tadpoles play an important role in (some) orientifold models with broken supersymmetry recently constructed [3, 4, 5].

The purpose of this letter is to explicitly find the background of the nonsupersymmetric tachyon-free strings in 10D: the type I model in ten dimensions [3], containing 32 $D\bar{9}$ antibranes and 32 $O9_-$ planes, and the heterotic $SO(16) \times SO(16)$ [7, 8]. For the two theories, there is no background with maximal $SO(10)$ Lorentz symmetry, a result to be expected by various considerations. We find explicitly the classical backgrounds with a 9D Poincaré symmetry. We find an unique solution for the Type I model and two independent solutions for the heterotic one. A remarkable feature (in the Type I and one of the two heterotic backgrounds) of the static solutions is that the tenth coordinate is dynamically compactified in the classical background. Furthermore, the effective nine-dimensional Planck and Yang-Mills constants are finite indicating that the low energy physics is nine-dimensional. Another classical background with maximal symmetry is a cosmological-type solution which we explicitly exhibit for the two theories. They both

have big-bang type curvature singularities.

In section 2 we briefly review the construction of the type I $USp(32)$ string. In section 3, we determine its classical background with maximal symmetry. In section 4, we consider the $SO(16) \times SO(16)$ heterotic string and finally we end in section 5 with a discussion of the solutions.

2 The Type I nonsupersymmetric $USp(32)$ string

The unique supersymmetric Type I model in ten dimensions is based on the gauge group $SO(32)$ and contains, by using a modern language (see, for example, [9]) 32 D9 branes and 32 $O9_+$ planes. There is however another, nonsupersymmetric tachyon-free model, with the same closed string spectrum at tree-level, containing 32 $\bar{D}9$ branes (i.e. branes of positive tension and negative RR charge) and 32 $O9_-$ planes (i.e. nondynamical objects with positive tension and positive RR charge).

The open string partition functions are [3]

$$\begin{aligned}\mathcal{K} &= \frac{1}{2(4\pi^2\alpha')^5} \int_0^\infty \frac{d\tau_2}{\tau_2^6} (V_8 - S_8) \frac{1}{\eta^8} , \\ \mathcal{A} &= \frac{N^2}{2} \frac{1}{(8\pi^2\alpha')^5} \int_0^\infty \frac{dt}{t^6} (V_8 - S_8) \frac{1}{\eta^8} , \\ \mathcal{M} &= \frac{N}{2} \frac{1}{(8\pi^2\alpha')^5} \int_0^\infty \frac{dt}{t^6} (V_8 + S_8) \frac{1}{\eta^8} ,\end{aligned}\tag{2.1}$$

where $\alpha' \equiv M_s^{-2}$ is the string tension and

$$V_8 = \frac{\theta_3^4 - \theta_4^4}{2\eta^4} , \quad S_8 = \frac{\theta_2^4 - \theta_1^4}{2\eta^4} ,\tag{2.2}$$

where θ_i are Jacobi functions and η the Dedekind function (see, for example, [9]). In (2.1) the various modular functions are defined on the double covering torus of the corresponding (Klein, annulus, Möbius) surface of modular parameter

$$\tau = 2i\tau_2 \quad (\text{Klein}) , \quad \tau = \frac{it}{2} \quad (\text{annulus}) , \quad \tau = \frac{it}{2} + \frac{1}{2} \quad (\text{Mobius}),\tag{2.3}$$

where $\tau = \tau_1 + i\tau_2$ is the modular parameter of the torus amplitude and t is the (one-loop) open string modulus.

As usual, the annulus and Möbius amplitudes have the dual interpretation of one-loop open string amplitudes and tree-level closed string propagation with the modulus l , related to the open string channel moduli by

$$\mathcal{K} : l = \frac{1}{2\tau_2} , \quad \mathcal{A} : l = \frac{2}{t} , \quad \mathcal{M} : l = \frac{1}{2t} . \quad (2.4)$$

From the closed string propagation viewpoint, V_8 describe the NS-NS sector (more precisely, the dilaton) and S_8 the RR sector, corresponding to an unphysical 10-form. The tadpole conditions can be derived from the $t \rightarrow 0$ ($l \rightarrow \infty$) limit of the amplitudes above and read

$$\mathcal{K} + \mathcal{A} + \mathcal{M} = \frac{1}{2} \frac{1}{(8\pi^2\alpha')^5} \int_0^\infty dl \{ (N+32)^2 \times 1 - (N-32)^2 \times 1 \} + \dots . \quad (2.5)$$

It is therefore clear that we can set to zero the RR tadpole by choosing $N = 32$, but we are forced to live with a dilaton tadpole. The resulting open spectrum is nonsupersymmetric (the closed spectrum is supersymmetric and given by the Type I supergravity) and contains the vectors of the gauge group $USp(32)$ and a fermion in the antisymmetric (reducible) representation. However, the spectrum is free of gauge and gravitational anomalies and therefore the model should be consistent. It is easy to realize from (2.1) that the model contains 32 $\bar{D}9$ branes and 32 $O9_-$ planes, such that the total RR charge is zero but NS-NS tadpoles are present, signaling breaking of supersymmetry in the open sector. The effective action, identified by writing the amplitudes (2.1) in the tree-level closed channel, contains here the bosonic terms

$$S = \frac{M_s^8}{2} \int d^{10}x \sqrt{-G} e^{-2\Phi} [R + 4(\partial\Phi)^2] - T_9 \int d^{10}x [(N+32)\sqrt{-G}e^{-\Phi} - (N-32)A_{10}] + \dots , \quad (2.6)$$

where T_9 is the D9 brane tension and we set to zero the RR two-form and the gauge fields, which will play no role in our paper. Notice in (2.6) the peculiar couplings of the dilaton

and the 10-form to antibranes and O9_− planes, in agreement with the general properties displayed earlier. The RR tadpole $N = 32$ is found in (2.6) simply as the classical field equation for the unphysical 10-form A_{10} .

The difference between the supersymmetric $SO(32)$ and nonsupersymmetric $USp(32)$ model described previously is in the Möbius amplitude describing propagation between (anti)branes and orientifold planes. Indeed, in the nonsupersymmetric case there is a sign change in the vector (or NS-NS in the closed channel) character V_8 . Both supersymmetric and nonsupersymmetric possibilities are however consistent with the particle interpretation and factorization of the amplitudes.

As noticed before, the NS-NS tadpoles generate scalar potentials for the corresponding (closed-string) fields, in our case the (10d) dilaton. The dilaton potential read

$$V \sim (N + 32)e^{-\Phi} \quad , \quad (2.7)$$

and in the Einstein basis is proportional to $(N + 32)\exp(3\Phi/2)$. It has therefore the (usual) runaway behaviour towards zero string coupling, a feature which is of course true in any perturbative construction. The dilaton tadpole means that the classical background, around which we must consistently quantize the string, cannot be the ten dimensional Minkowski vacuum and solutions with lower symmetry must be searched for. Once the background is correctly identified, there is no NS-NS tadpole anymore, of course.

3 The classical background of the nonsupersymmetric $USp(32)$

Type I string

We are searching for classical solutions of the effective lagrangian (2.6) of the model (2.1) in the Einstein frame, which reads

$$S_E = \frac{1}{2k^2} \int d^{10}x \sqrt{-G} [R - \frac{1}{2}(\partial\Phi)^2] - T_9^E \int d^{10}x [(N + 32)\sqrt{-G}e^{\frac{3\Phi}{2}} - (N - 32)A_{10}] + \dots \quad , \quad (3.1)$$

where T_9^E indicate that the tension here is in the Einstein frame. The maximal possible symmetry of the background of the model described in the previous paragraph has a nine dimensional Poincaré isometry and is of the following form

$$ds^2 = e^{2A(y)} \eta_{\mu\nu} dx^\mu dx^\nu + e^{2B(y)} dy^2, \quad \Phi = \Phi(y), \quad (3.2)$$

where $\mu, \nu = 0 \cdots 8$ and the antisymmetric tensor field from the RR sector, the gauge fields and all fermion fields are set to zero. The Einstein and the dilaton field equations with this ansatz are given by

$$\begin{aligned} 36(A')^2 + 8A'' - 8A'B' + \frac{1}{4}(\Phi')^2 &= -\alpha_E e^{2B+3\Phi/2}, \\ 36(A')^2 - \frac{1}{4}(\Phi')^2 &= -\alpha_E e^{2B+3\Phi/2}, \\ \Phi'' + (9A' - B')\Phi' &= 3\alpha_E e^{2B+3\Phi/2}, \end{aligned} \quad (3.3)$$

where we defined $\alpha_E = (N + 32)k^2 T_9^E = 64k^2 T_9^E$ and $A' \equiv dA/dy$, etc. The function B can be gauge-fixed by using the reparametrisation invariance of the above equations. It is convenient to choose the coordinate y where $B = -3\Phi/4$ so that the exponential factors in the equations (3.3) disappear. In this coordinate system, the second equation in (3.3) is solved in terms of one function¹ f

$$A' = \frac{1}{6}\sqrt{\alpha_E} \, shf, \quad \Phi' = 2\sqrt{\alpha_E} \, chf. \quad (3.4)$$

The two other field equations become then

$$\begin{aligned} \frac{4}{3}\sqrt{\alpha_E} \, f' chf + \alpha_E e^{2f} &= -\alpha_E, \\ 2\sqrt{\alpha_E} \, f' shf + \frac{3}{2}\alpha_E e^{2f} &= \frac{3}{2}\alpha_E, \end{aligned} \quad (3.5)$$

and the solution is then

$$e^{-f} = \frac{3}{2}\sqrt{\alpha_E} \, y + c, \quad (3.6)$$

¹It can be checked that the other possible sign choices in (3.4) lead to the same solution.

where c is a constant. By a choice of the y origin and rescaling of the x^μ coordinates, the final solution in the Einstein frame reads

$$\begin{aligned}\Phi &= \frac{3}{4}\alpha_E y^2 + \frac{2}{3}\ln|\sqrt{\alpha_E}y| + \Phi_0, \\ ds_E^2 &= |\sqrt{\alpha_E}y|^{1/9}e^{-\alpha_E y^2/8}\eta_{\mu\nu}dx^\mu dx^\nu + |\sqrt{\alpha_E}y|^{-1}e^{-3\Phi_0/2}e^{-9\alpha_E y^2/8}dy^2.\end{aligned}\quad (3.7)$$

For physical purposes it is also useful to display the solution in the string frame, related as usual by a Weyl rescaling $G \rightarrow e^{\frac{\Phi}{2}}G$ to the Einstein frame

$$A_s = A + \frac{1}{4}\Phi, \quad B_s = B + \frac{1}{4}\Phi. \quad (3.8)$$

In the string frame the solution reads

$$\begin{aligned}g_s &\equiv e^\Phi = e^{\Phi_0}|\sqrt{\alpha}y|^{2/3}e^{3\alpha y^2/4}, \\ ds^2 &= |\sqrt{\alpha}y|^{4/9}e^{\Phi_0/2}e^{\alpha y^2/4}\eta_{\mu\nu}dx^\mu dx^\nu + |\sqrt{\alpha}y|^{-2/3}e^{-\Phi_0}e^{-3\alpha y^2/4}dy^2,\end{aligned}\quad (3.9)$$

where $\alpha = 64M_s^{-8}T_9$. The solution (3.9) displays two timelike singularities, one at the origin $y = 0$ and one at infinity $y = +\infty$, so that the range of the y coordinate is $0 < y < +\infty$. The dilaton, on the other hand, vanishes at $y = 0$ and diverges at $y = +\infty$. The brane solution found above (3.9) has a striking feature. Suppose the y coordinate is noncompact $0 < y < +\infty$. In curved space however, the real radius R_c is given by the integral

$$2\pi R_c = \int_0^\infty dy e^B = e^{-\Phi_0/2}\alpha^{-\frac{1}{2}} \int_0^\infty \frac{du}{u^{1/3}} e^{-3u^2/8}, \quad (3.10)$$

where $u = \sqrt{\alpha}y$. The result is finite, meaning that despite appearances the tenth coordinate is actually compact. The topology of the solution is thus a ten dimensional manifold with two boundaries at $y = 0$ and $y = +\infty$ that is $R^9 \times S^1/Z_2$. Moreover, it can be argued that gravity and gauge fields of the D $\bar{9}$ branes are confined to the nine-dimensional non-compact subspace, by computing the nine-dimensional Planck mass and gauge couplings, respectively

$$\begin{aligned}M_P^7 &= M_s^8 \int_0^\infty dy e^{7A_s+B_s-2\Phi} = M_s^8 \alpha^{-\frac{1}{2}} e^{-3\Phi_0/4} \int_0^\infty \frac{du}{u^{1/9}} e^{-3u^2/4}, \\ \frac{1}{g_{YM}^2} &= M_s^6 \int_0^\infty dy e^{5A_s+B_s-\Phi} = M_s^6 \alpha^{-\frac{1}{2}} e^{-\Phi_0/4} \int_0^\infty du u^{1/9} e^{-u^2/2}.\end{aligned}\quad (3.11)$$

Both of them are finite, which indeed suggest that gravity and gauge fields of the $D\bar{9}$ branes are confined to the nine-dimensional subspace. The relations (3.10) and (3.11) are in sharp contrast with the usual flat space relations obtained by compactifying the ten-dimensional theory down to nine-dimensions on a circle of radius R_c

$$M_P^7 \sim e^{-2\Phi_0} R_c M_s^8, \quad \frac{1}{g_{YM}^2} \sim e^{-\Phi_0} R_c M_s^6. \quad (3.12)$$

By using exactly the same method we find a cosmological solution for the $USp(32)$ non-supersymmetric Type I model by searching a homogeneous metric of the form

$$ds^2 = -e^{2B(t)} dt^2 + e^{2A(t)} \delta_{\mu\nu} dx^\mu dx^\nu, \quad \Phi = \Phi(t), \quad (3.13)$$

where t is a time coordinate. The solution is easily found by following the steps which led to (3.7) and (3.9). The result in the string frame is

$$\begin{aligned} g_s &= e^\Phi = e^{\Phi_0} |\sqrt{\alpha} t|^{2/3} e^{-3\alpha t^2/4}, \\ ds^2 &= -|\sqrt{\alpha} t|^{-2/3} e^{-\Phi_0} e^{3\alpha t^2/4} dt^2 + |\sqrt{\alpha} t|^{4/9} e^{\Phi_0/2} e^{-\alpha t^2/4} \delta_{\mu\nu} dx^\mu dx^\nu. \end{aligned} \quad (3.14)$$

The metric has a spacelike curvature singularities at $t = 0$ and $t = +\infty$. The laps separating these two singularities, to be interpreted as the real time parameter,

$$\tau = \int dt |\sqrt{\alpha} t|^{-1/3} e^{-\Phi_0/2} e^{3\alpha t^2/8}, \quad (3.15)$$

is infinite.

4 Classical background of the $SO(16) \times SO(16)$ heterotic string

In ten dimensions there is a unique tachyon-free non-supersymmetric heterotic string model [7, 8]. It can be obtained from the two supersymmetric heterotic strings as a Z_2 orbifold. The resulting bosonic spectrum comprises the gravity multiplet: graviton, dilaton, antisymmetric tensor and gauge bosons of the gauge group $SO(16) \times SO(16)$. Compactifications of this theory were considered in [10], and its strong coupling behavior in nine

dimensions was examined in [11]. It has been shown [8] that the cosmological constant (the partition function on the torus) at one loop is finite and positive, furthermore its approximate value is given by

$$\Lambda \approx M_s^{10} \frac{2^6}{(2\pi)^{10}} \times 5.67 . \quad (4.1)$$

The effective low energy action for the gravity multiplet is the same as before except for the cosmological constant term, which now reads

$$- \Lambda \int \sqrt{-G} \quad (4.2)$$

in the string metric. The absence of the dilaton in this term reflects the one-loop nature of the cosmological constant². The same ansatz of the Einstein metric as in the previous paragraph leads, in the Einstein frame, to the equations

$$\begin{aligned} 36(A')^2 + 8A'' - 8A'B' + \frac{1}{4}(\Phi')^2 &= -\beta_E e^{2B+5\Phi/2} , \\ 36(A')^2 - \frac{1}{4}(\Phi')^2 &= -\beta_E e^{2B+5\Phi/2} , \\ \Phi'' + (9A' - B')\Phi' &= 5\beta_E e^{2B+5\Phi/2} , \end{aligned} \quad (4.3)$$

where we defined $\beta_E = \Lambda_E k^2$. The gauge which eliminates the exponential factors is now $B = -5\Phi/4$. After solving the second equation in (4.3) $A' = \sqrt{\beta_E} sh(h)/6$, $\Phi' = 2\sqrt{\beta_E} ch(h)$, the remaining equations give

$$e^h = \frac{1}{2} \frac{e^{\sqrt{\beta_E} y} + \epsilon e^{-\sqrt{\beta_E} y}}{e^{\sqrt{\beta_E} y} - \epsilon e^{-\sqrt{\beta_E} y}}, \quad (4.4)$$

where $\epsilon = \pm 1$. An important difference with respect to the type I solution is that here we have two non-equivalent (that is, not related by coordinate transformations) solutions corresponding to $\epsilon = 1$ or -1 . Let us first consider the $\epsilon = 1$ case. The solution in the Einstein frame reads

$$\Phi = \Phi_0 + \frac{1}{2} \ln |sh \sqrt{\beta_E} y| + 2 \ln (ch \sqrt{\beta_E} y) ,$$

²A direct computation shows also that the one-loop dilaton tadpole is non-zero and proportional to the cosmological constant Λ .

$$ds_E^2 = |sh(\sqrt{\beta_E y})|^{\frac{1}{12}} \left(ch(\sqrt{\beta_E y}) \right)^{-\frac{1}{3}} dx^2 + e^{-5\Phi_0/2} |sh(\sqrt{\beta_E y})|^{\frac{5}{4}} \left(ch(\sqrt{\beta_E y}) \right)^{-5} dy^2 \quad (4.5)$$

and in the string frame

$$e^{2\Phi} = e^{2\Phi_0} |sh(\sqrt{\beta} y)| \left(ch(\sqrt{\beta} y) \right)^4 ,$$

$$ds^2 = e^{\Phi_0/2} |sh(\sqrt{\beta} y)|^{\frac{1}{3}} \left(ch(\sqrt{\beta} y) \right)^{\frac{2}{3}} dx^2 + e^{-2\Phi_0} |sh(\sqrt{\beta} y)|^{-1} \left(ch(\sqrt{\beta} y) \right)^{-4} dy^2 , \quad (4.6)$$

where here $\beta = \Lambda M_s^{-8}$. In both metrics, the solution has two timelike singularities at $y = 0$ and $y = \infty$. These singularities are separated by a finite distance which in the string frame reads

$$2\pi R_c = (\beta)^{-1/2} e^{-\Phi_0} \int_0^{+\infty} du (sh u)^{-1/2} (ch u)^{-2} . \quad (4.7)$$

The spacetime has therefore the topology of a nine dimensional Minkowski space times an interval. Notice that the nine dimensional Planck mass

$$M_p^7 = M_s^8 (\beta)^{-1/2} e^{-5\Phi_0/4} \int_0^\infty du (sh u)^{-1/3} (ch u)^{-11/3} , \quad (4.8)$$

as well as the 9D Yang-Mills coupling

$$\frac{1}{g_{YM}^2} = M_s^6 e^{-7\Phi_0/4} (\beta)^{-1/2} \int_0^\infty du (sh u)^{-2/3} (ch u)^{-13/3} , \quad (4.9)$$

are finite. This fact together with the finitude of the length of the tenth coordinate indicate that the low energy processes are described by a 9D theory. The coupling constant vanishes at 0 and becomes infinite at large y . We shall comment more on this fact in the conclusion.

The second solution corresponding to $\epsilon = -1$ can be obtained by exchanging $(ch(\sqrt{\beta} y))$ with $|sh(\sqrt{\beta} y)|$ in the solutions (4.5) and (4.6). This solution has two singularities at 0 and ∞ , however the nine-dimensional Planck and Yang-Mills constants as well as the length of the tenth coordinate are infinite.

The cosmological solution invariant with respect to the nine dimensional Euclidian group can be also readily found and it reads in the string frame

$$ds^2 = -e^{-2\Phi_0} (\sin \sqrt{\beta} t)^{-1} (\cos \sqrt{\beta} t)^{-4} dt^2 + e^{\Phi_0/2} (\sin \sqrt{\beta} t)^{1/3} (\cos \sqrt{\beta} t)^{2/3} dx^2 ,$$

$$e^{2\Phi} = e^{2\Phi_0} (\sin \sqrt{\beta} t) (\cos \sqrt{\beta} t)^4. \quad (4.10)$$

The variable t in these equations belongs to the interval $[0, \pi/(2\sqrt{\beta})]$. At the boundaries in $t = 0$ and $t = \pi/(2\sqrt{\beta})$ the metric develops curvature singularities. The time separating these two singularities is infinite:

$$\tau = \int_0^{\pi/(2\sqrt{\beta})} dt (\sin \sqrt{\beta} t)^{-1/2} (\cos \sqrt{\beta} t)^{-2} = \infty. \quad (4.11)$$

Notice that the solution obtained by exchanging the sine and cosine in the above equations is not a new one since it can be obtained by shifting the time coordinate $t \rightarrow \pi/(2\sqrt{\beta}) - t$.

5 Discussion

Before discussing the solutions we have found we should mention that there exists another interesting non supersymmetric and tachyon-free model in ten dimensions [12]: it is an orientifold of the type 0B string with a projection that removes the tachyon and introduces an open sector with the gauge group $U(32)$. This model has both a one loop (positive) cosmological constant and a disk dilaton tadpole so we have a sum of two terms in the low energy effective action

$$- \Lambda_1 \int \sqrt{-G} e^{-\Phi} - \Lambda_2 \int \sqrt{-G}, \quad (5.1)$$

in the string metric. The first term is of the type we encountered in the $USp(32)$ case and the second is analogous to the one we met in the $SO(16) \times SO(16)$ case. There is no simple gauge choice for B that renders the equations as simple as before. However qualitatively the solution should behave as the $USp(32)$ case for small y where the coupling constant is small and the behavior for large y should resemble that of the $SO(16) \times SO(16)$ case. In particular we expect the radius of the tenth dimension to be compactified (since the divergence of the radius in the second solution of $SO(16) \times SO(16)$ is due to the behavior at the origin).

We have determined the maximally symmetric solutions to the low energy equations of two tachyon-free non-supersymmetric strings. To what extent can we consider these solutions as representing the vacuum of these string theories ? Perturbatively, there are two kind of string corrections to the low energy effective action. The first ones are α' corrections involving the string oscillators and the second ones are g_s string loop corrections. A common feature of the solutions we found is that the conformal factor e^{2A} (in nine dimensions) as well as the string coupling vanish at the origin and diverge at $y = \infty$. The effective string scale at coordinate y being given by $M_s^2(y) = M_s^2 e^{2A_s}$, we expect α' corrections to be important at the origin and the loop corrections to be dominant at infinity. So strictly speaking we cannot trust the classical solution near the two singularities where interesting string physics would occur. Another less ambitious question, is the classical stability of our solutions. That is, do small perturbations around the background we found destroy the solution ? The answer to this question is closely related to the determination of the Kaluza-Klein excitations [13].

Another common feature of the static solution of the Type I model (3.9) and of the first heterotic solution (4.6) is that the effective Yang-Mills and Planck constants are finite which means that the gravitational and gauge physics is effectively nine-dimensional. A remarkable feature of the static solutions in the low energy approximation is the spontaneous compactification of one coordinate. Whether this feature will survive the string and loop corrections is an interesting and open question. The fact that the nine dimensional metric is flat in spite of the ten-dimensional cosmological constant is in the spirit of the higher dimensional mechanisms which try to explain the vanishing of the (effective) cosmological constant [14]. As in the new approaches to this problem discussed in [14], however, a better understanding of the naked singularities present in our solutions is needed in order to substantiate this claim.

A notable difference between the type I and heterotic solutions is that the type I back-

ground is unique, whereas there are two classical heterotic backgrounds with physically very different properties. It is possible that this is due to the low energy approximation and that this degeneracy will be lifted by string corrections.

All of the nontrivial features of the solutions are due to the presence of dilaton tadpoles. The latter are generic for non-supersymmetric string models. According to the Fischler-Susskind mechanism, the quantization around the classical solution should lead to finite string amplitudes [1, 9]. It would be interesting to confirm this explicitly for the present models.

Acknowledgments

We are grateful to C. Angelantonj and A. Sagnotti for illuminating discussions on nonsupersymmetric strings and to C. Grojean and S. Lavignac for discussions concerning naked singularities in relation with the proposals [14]. E.D. would like to thank the Theory Group at LBNL–Berkeley for warm hospitality during the final stage of this work.

References

- [1] W. Fischler and L. Susskind, *Phys. Lett.* **B171** (1986) 383; *Phys. Lett.* **B173** (1986) 262.
- [2] J. Polchinski and Y. Cai, *Nucl. Phys.* **B296** (1988) 91; J. Polchinski, *Phys. Rev. Lett.* **75** (1995) 4724.
- [3] S. Sugimoto, hep-th/9905159.
- [4] I. Antoniadis, E. Dudas and A. Sagnotti, *Phys. Lett.* **B464** (1999) 38; G. Aldazabal and A.M. Uranga, hep-th/9908072; G. Aldazabal, L.E. Ibánñez and F. Quevedo, hep-th/9909172 and hep-ph/0001083; C. Angelantonj, I. Antoniadis, G. D’Appollonio, E. Dudas and A. Sagnotti, hep-th/9911081.

- [5] C. Bachas, hep-th/9503030.
- [6] A. Sen, *J. High Energy Phys.* **9806** (1998) 007; *J. High Energy Phys.* **9808** (1998) 010; *J. High Energy Phys.* **9809** (1998) 023; *J. High Energy Phys.* **9812** (1998) 021.
- [7] L. Dixon and J. Harvey, *Nucl. Phys.* **B274** (1986) 93.
- [8] L. Alvarez-Gaume, P. Ginsparg, G. Moore and C. Vafa, *Phys. Lett.* **B171** (1986) 155.
- [9] J. Polchinski, String theory, Cambridge Univ. Press, 1998.
- [10] V.P. Nair, A. Shapere, A. Strominger and F. Wilczek, *Nucl. Phys.* **B287** (1987) 402.
- [11] J.D. Blum and K.R. Dienes, hep-th/9707160; hep-th/9707148.
- [12] A. Sagnotti, hep-th/9509080; hep-th/9702093.
- [13] E. Dudas, J. Mourad, work in progress.
- [14] V.A. Rubakov and M.E. Shaposhnikov, *Phys. Lett.* **B125** (1983) 139; J. de Boer, E. Verlinde and H. Verlinde, hep-th/9912012; E. Verlinde and H. Verlinde, hep-th/9912018; N. Arkani-Hamed et al., hep-th/0001197; S. Kachru, M. Schulz and E. Silverstein, hep-th/0001206; D. Youm, hep-th/0002147; S. Forste et al., hep-th/0002164; C. Csaki et al., hep-th/0004133.